hep-th/9708049
UTTG–22–97

Phenomenological Aspects of F-theory⋆

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ABSTRACT

Stabilizing a heterotic string vacuum with a large expectation value of the dilaton and simultaneously breaking low-energy supersymmetry is a long-standing problem of string phenomenology. We reconsider these issues in light of the recent developments in F–theory.

⋆ Research supported in part by: the NSF, under grant PHY–95–11632 (V. K.); the Robert A. Welch Foundation (V. K.); the German–Israeli Foundation for Scientific Research (J. L.); the NATO, under grant CRG 931380 (both authors).

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Since the inception of the heterotic string theory, a very large number of phenomenological string models were constructed based upon various candidate ground states of the string. Although such models differ from each other in countless details, certain ground rules have been firmly established as either universal consequences of the perturbative heterotic string theory or else essential for obtaining the correct Standard Model phenomenology.[1] Generally, a string model has four spacetime dimensions, $\mathcal{N}=1$ supersymmetry and a large gauge symmetry $G=\prod_a G_a$ including the $SU(3)\times SU(2)\times U(1)$ of the Standard Model as well as additional, 'hidden' factors. At the tree level, all the gauge couplings g_a are controlled by the expectation value of the dilaton field S,

$$\frac{4\pi}{g_a^2} \equiv \frac{1}{\alpha_a} \;=\; k_a \,\langle \operatorname{Re} S\rangle \tag{1}$$

(k_a being fixed integer or rational coefficients). The universality of this relation naturally leads to the desired GUT-like pattern of the Standard Model's gauge couplings.* The perturbative string theory suffers from an exact degeneracy which leaves $\langle S\rangle$ completely undetermined; likewise, the vacuum expectation values of several other moduli fields (collectively denoted T) are also indeterminate to all orders of the string perturbation theory.

The lifting of this degeneracy as well as spontaneous supersymmetry breakdown can be accomplished with the help of field-theoretical non-perturbative effects arising from asymptotically free (and hence infrared-strong) hidden sectors of the gauge group.[4] In the simplest scenario, a confining hidden sector generates a dynamical superpotential $W\sim\Lambda_{\rm hid}^3$ where

$$\Lambda_{\rm hid} \;\sim\; e^{-2\pi kS/b}\, M_{\rm Pl} \tag{2}$$

is the confinement scale and b the appropriate β-function coefficient. Taking several

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$\star$ The string analogue of the GUT scale $3\cdot 10^{17}$ GeV does not exactly coincide with the phenomenologically favorite value $M_{\rm GUT}\approx 2\cdot 10^{16}$ GeV, but the discrepancy is small enough to be explainable (in principle) in terms of the perturbative string threshold corrections.[2] It has also been suggested that this mismatch can disappear in strongly coupled heterotic vacua.[3]

such hidden sectors together and allowing for moduli-dependent pre-exponential factors, one generally has[5]

$$W_{\text{eff}}(S,T) \;=\; M_{\text{Pl}}^3 \sum_a C_a(T)\, e^{-6\pi k_a S/b_a} \tag{3}$$

(a runs over the confining hidden sectors), which leads to an effective scalar potential $V(S,T) = e^K \left[|DW|^2 - 3\,|W|^2\right]$. Phenomenologically, this effective potential should have a stable minimum with spontaneously broken supersymmetry and zero cosmological constant. Furthermore, the observable sector (*i.e.*, the Standard Model) should feel the breakdown of supersymmetry at the electroweak scale M_W; this requires $W_{\text{eff}} = O(M_W M_{\text{Pl}}^2)$ or equivalently confinement scales Λ_{hid} in the 10^{13} GeV to 10^{14} GeV range.

In all other scenarios, the hierarchy $M_W \ll M_{\text{Pl}}$ also follows from $\Lambda_{\text{hid}} \ll M_{\text{Pl}}$ for some kind of a hidden sector. According to eq. (2), this requires a rather large expectation value of the dilaton field, typically $\langle ReS\rangle \gtrsim 10$ or more. Likewise, extrapolating the Supersymmetric Standard Model all the way up to the GUT scale and using eq. (1), one needs $\langle ReS\rangle \approx \alpha_{\text{GUT}}^{-1} \approx 23$.[6] Unfortunately, it is extremely difficult to stabilize the dilaton at such a large value using only field-theoretical non-perturbative effects. According to Dine and Seiberg,[7] for any string model with unbroken supersymmetry at the tree level, the effective potential exponentially asymptotes to zero in the weak coupling regime $\operatorname{Re} S \to \infty$ and hence, the stable minima of the potential, if any, must lie at strong coupling.† For example, the superpotential (3) with a generic Kähler function $K(S,T)$ and no special tuning of the coefficients $C_a(T)$ and k_a/b_a, leads to stable vacua only when some of the exponential factors $e^{-6\pi k_a S/b_a}$ are not too small ($O(1)$) and hence

$$\langle \operatorname{Re} S\rangle \;\lesssim\; O(1/6\pi)\, \max_a (b_a/k_a)\,. \tag{4}$$

An example of a hidden sector with a large b/k ratio is the unbroken second E_8 gauge group of the heterotic string ($b = 90$, $k = 1$); depending on how seriously

† See also the discussion in refs. [8].

one takes the numerical factors in eq. (4), it might be barely consistent with stable $\langle ReS\rangle \approx 23$. At the same time however, one would have $\Lambda_{E_8} = O(M_{\rm Pl})$, which leads to spontaneous supersymmetry breakdown close to the Planck scale and hence no hierarchy.

From the purely field-theoretical point of view, one might attempt to solve the problem by employing several large hidden gauge factors[5] to stabilize a large $\langle \mathrm{Re}\, S\rangle$ without breaking supersymmetry and then add yet another hidden sector with $\Lambda \ll M_{\rm Pl}$ for the express purpose of breaking supersymmetry at a hierarchically small scale.‡ However, from the heterotic string's point of view, this scenario — or any other scenario which needs very large or complicated hidden sectors — conflicts with the universal central charge constraint, which limits the rank of the entire (perturbative) four-dimensional gauge group:

$$\mathrm{rank}(G) \leq 22\,; \tag{5}$$

this leaves rather limited room for the hidden sectors. Consequently, the perturbative heterotic string theory with only field-theoretical non-perturbative corrections has extreme difficulty combining a stable vacuum with a large dilaton expectation value and a large hierarchy.

The inherently stringy non-perturbative effects are now gradually becoming understood in terms of duality relations between various string theories, M–theory and F–theory. In particular, the $\mathcal{N} = 1$, $d = 4$ compactifications of the heterotic string are dual to F–theory compactifications on elliptically fibered Calabi-Yau fourfolds.[9,10,11,12] This duality shows that the perturbative heterotic string theory often reveals only a small part of the ultimate four-dimensional gauge group G

‡ The Krasnikov mechanism — stabilizing a large $\langle \mathrm{Re}\, S\rangle$ by using two or more hidden gauge factors — can be fine-tuned to work with modestly sized hidden gauge group provided their β-functions are very close but not exactly equal, $b_1 \approx b_2 \gg |b_1 - b_2| > 0$. With the help of the moduli fields T — and even more fine tuning — one may achieve a hierarchical supersymmetry breakdown. However, the extreme fine tuning required by this scenario makes it rather marginal for model-building purposes.

while many additional gauge fields arise from singularities of the heterotic compactification where the perturbation theory breaks down.[13] § Only the perturbative gauge couplings g_a are governed by the dilaton S; the non-perturbative couplings are dilaton-independent and are instead controlled by some combinations of the moduli fields T.[16] From the F–theory point of view, however, all the gauge groups G_a have equal status and S is no different from the other moduli fields. Only in the corner of the F-theory moduli space which is dual to the weakly coupled heterotic string does the S field acquire its special properties. There are no known constraints on the size or the variety of the non-perturbative gauge symmetries.[17] The current record holder among the F–theory compactifications has 251 simple gauge group factors of total non-abelian rank of 302896, the biggest factors being $SO(7232)$ and $Sp(3528)$.[12]

What then are the implications of this non-perturbative bounty of gauge fields for the string model building? On one hand, we no longer have any general constraints — or guidelines — for the hidden sectors of string models. On the other hand, it is precisely the absence of constraints such as (5) that makes it easy to obtain a small $\alpha_{\rm GUT}$ in a stable vacuum. For example, imagine a model where the same combination $\tilde{S}$ of moduli fields governs the gauge couplings of the Standard Model and also of several *large* confining hidden gauge groups. In this model, the effective superpotential for $\tilde{S}$ and the other moduli looks exactly like (3) (modulo replacement $S \to \tilde{S}$) and hence the minima of the resulting effective potential are generally found at $6\pi \langle \mathrm{Re}\, \tilde{S} \rangle = O(b_{\rm hid})$. This is compatible with $\mathrm{Re}\, \tilde{S} \approx 23$ for $b_{\rm hid} \sim 400$ — as in *e.g.*, pure-gauge $SO(140)$ — which would be easily obtainable in F–theory compactifications (but quite out of reach of the perturbative heterotic theory).

In a more realistic model, different gauge couplings $\alpha_a(T)$ may be controlled by different dilaton-like combinations of the moduli fields T (which by abuse of

§ This is distinct from the perturbation theory breakdown due to a large overall ten-dimensional heterotic string coupling. That regime is best described in terms of the dual M–theory[14]; some of its phenomenological implications are discussed in refs. [3,15].

notations now include the heterotic S field as well). In this case, one has

$$W_{\text{eff}}(T) \;=\; M_{\text{Pl}}^{3}\sum_{a} C_{a}(T)\, e^{-6\pi/b_a\alpha_a(T)} \tag{6}$$

or even a more complicated non-linear combination of the $e^{-6\pi/b_a\alpha_a(T)}$ if there are hidden matter fields that are charged under several hidden gauge groups at once. Generally, there are also inherently F–theoretical instantonic contributions, although they are believed to be smaller than those of confining hidden sectors.[10] The precise behavior of the resulting scalar potential can only be analyzed on the model-by-model basis, but a crude order-of-magnitude analysis suggests that its stable minima (if any) should have

$$\frac{6\pi}{\alpha_{\text{hid}}(\langle T\rangle)} \;=\; O(b_{\text{hid}}). \tag{7}$$

Again, we see that large hidden sectors naturally lead to small α_{hid}.

In a generic F–theory compactification, however, small α_{hid} do not necessarily imply a weakly coupled Standard Model. Furthermore, there is no longer an automatic GUT-like unification of the Standard Model's couplings themselves. Both of these features — which perturbatively followed from eq. (1) — now have to be imposed as phenomenological constraints on the F-theory models. Specifically, the three Standard Model's couplings should be governed by the same modulus,* which is also involved with the hidden gauge couplings (7) and thus obtains a large expectation value.

Notice that eq. (7) implies $\Lambda_{\text{hid}} \sim M_{\text{Pl}}$ and hence $W_{\text{eff}} \sim M_{\text{Pl}}^{3}$. Therefore, it is imperative that the resulting effective potential does *not* lead to spontaneous supersymmetry breakdown — otherwise, supersymmetry would be broken right at the Planck scale for all sectors of the theory and there would be no hierarchy. Likewise, we do not want a Planck-scale cosmological constant. This gives us two

$\star$ This suggest that in F–theory one should seriously consider a possibility of an actual field-theoretical Grand Unification of the $SU(3)\times SU(2)\times U(1)$ into a simple gauge group with a single $\alpha_{\text{GUT}}(T)$.

more phenomenological requirements that any viable F–theory model must satisfy. Mathematically, these requirements amount to a constraint on the holomorphic $W_{\rm eff}$,† which (in principle) allows one to decide the viability of any particular model in terms of exactly computable quantities.

Having survived the Planck-scale physics unbroken, supersymmetry should be eventually broken down at a hierarchically lower scale. This can be accomplished by additional hidden sectors — fortunately, they are easily available in F–theory. Such a sector needs the following features: A week coupling $\alpha \sim \alpha_{\rm GUT}$, a modest amount of asymptotic freedom $b \sim 10$ — which together provide for the hierarchy $\Lambda \ll M_{\rm Pl}$ — and most importantly, supersymmetry-breaking infrared dynamics.[18] There are several possibilities for such supersymmetry breakdown and for the way it affects the Standard Model. In the simplest scenario, vacuum stabilization and supersymmetry breaking result from completely separate hidden sectors: Large gauge groups (*e.g.*, pure-gauge $SO(140)$) have $\Lambda \sim M_{\rm Pl}$ and create $W_{\rm eff}$ that stabilizes *all* the moduli T of the F–theory right at the Planck scale, while a sector such as $SU(5)$ with $\mathbf{10} + \bar{\mathbf{5}}$ matter has Λ in a multi-TeV range and breaks supersymmetry dynamically without any help from the moduli fields or supergravity.[18] Finally, an abelian hidden gauge field communicates the supersymmetry breaking to the Standard Model *à la* Dine-Nelson.[19] Alternatively, the feed-down of supersymmetry breaking to the Standard Model can proceed through the supergravity- or moduli-mediated interactions. Also some of the moduli fields can avoid getting Planck-scale masses and survive to participate in the supersymmetry-breaking process.‡ Unfortunately, any involvement of the moduli fields in supersymmetry breaking is likely to destroy the charge universality of the squark and slepton masses of

† Specifically, one needs a simultaneous solution of holomorphic eqs. $W_{\rm eff} = \partial_T W_{\rm eff} = 0$, which is automatically a local minimum of the scalar potential with unbroken supersymmetry and zero cosmological constant. Whether or not it is the global minimum of the potential depends on the Kähler function.

‡ In this scenario, $\Lambda_{\rm hid}$ has to be $10^{13} - 10^{14} GeV$ in order to generate an effective superpotential for the surviving moduli fields, and it is this superpotential that leads to the spontaneous supersymmetry breakdown. This is rather similar to the heterotic toy models in which a large $\langle S \rangle$ is fixed by hand, while the spontaneous supersymmetry breakdown is induced by the effective potential for the T modulus.[20]

the Supersymmetric Standard Model.[21] From this point of view, the Dine–Nelson scenario appears more attractive.

Let us now summarize the key points: The non-perturbative string theory or F–theory allow for essentially unilimited hidden sectors. This makes it relatively easy to arrange for a stable vacuum state where supersymmetry is broken at a hierarchically low scale. On the other hand, the GUT-like unification of the Standard Model's gauge couplings is no longer automatic but instead has to be imposed as a phenomenological constraint. We do not propose any specific models but merely outline a general scenario for obtaining viable phenomenology from the F–theory. Indeed, it is hard to be specific without a better understanding of the moduli dependence of the gauge couplings in F–theory or even general rules for obtaining the spectra of the charged matter fields. However, we believe our scenario is a useful starting point for future work.

Acknowledgements: We would like to thank the CERN theory group for hospitality and P. Mayr for helpful discussions.

The research of V. K. is supported in part by the NSF, under grant PHY–95–11632, and by the Robert A. Welch Foundation. J. L. is supported in part by the German–Israeli Foundation for Scientific Research. The collaboration of the two authors is additionally supported by the NATO, under grant CRG 931380.

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